Rare $B \to K^* \nu \bar{\nu}$ decay beyond standard model

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Abstract

Using the most general, model–independent form of effective Hamiltonian, the exclusive, rare $B \to K^* \nu \bar{\nu}$ decay is analyzed. The sensitivity of the branching ratios and missing mass–squared spectrum to the new Wilson coefficients is discussed.

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1 Introduction

Started operating, two B–factories BaBar and Belle [1] open an excited era for studying B meson physics and its rare decays. The main physics program of these factories constitutes a detailed study of CP violation in B_d meson and precise measurement of rare flavor–changing neutral current (FCNC) processes. The rare decays of B mesons take place via FCNC and appear in Standard Model (SM) at loop level. For this reason there appears a real possibility for checking the gauge structure of SM at loop level. On the other side rare decays are very sensitive to the new physics beyond SM and their study is hoped to shed light on the existence of new particles before they are produced at colliders.

As has already been mentioned, one main goal of the B physics program is to find inconsistencies within the SM, in particular to find indications for new physics in the flavor and CP violating sectors [2]. In general, new physics effects manifest themselves in rare B meson decays either through new contributions to the Wilson coefficients that exist in the SM or by introducing new structures in the effective Hamiltonian which are absent in the SM (see [3]–[12]). Moreover, one can add new CP–violating phases and modify the flavor changing neutral current.

Currently the main interest is focused on the rare B meson decays, for which SM predicts "large" branching ratios. The rare $B \to K^* \nu \bar{\nu}$ decay is such a decay that plays a special role, both from experimental and theoretical point of view.

At quark level the FCNC $b \to s \nu \bar{\nu}$ decay is described in framework of the SM by the effective Hamiltonian [13]

$$\mathcal{H}_{eff} \;=\; \frac{G_F \alpha}{2\sqrt{2}\pi \sin^2\theta_W} V_{tb} V_{ts}^* X(x) \bar{b} \gamma^\mu (1-\gamma_5) s \bar{\nu} \gamma_\mu (1-\gamma_5) \nu \ , \tag{1}$$

where G_F is the Fermi coupling constant, α is the fine structure coupling constant, θ_W is the Weinberg angle, V_{ij} is the Cabibbo–Kobayashi–Maskawa (CKM) matrix elements and

$$X(x) = X_0(x) + \frac{\alpha_s}{4\pi} X_1(x) \ , \tag{2}$$

where

$$X_0 = \frac{x}{8} \left[\frac{x+2}{x-1} + \frac{3(x-2)}{(x-1)^2} \ln x \right] \ , \tag{3}$$

and $x = m_t^2/m_W^2$. Explicit form of $X_0(x)$ and $X_1(x)$ can be found in [14] and [13], respectively. Note that $X_1(x)$ gives about 3% contribution to the $X_0(x)$ term. The main attractive property of (1) is that the $b \to s \nu \bar{\nu}$ decay is governed only by a single operator and is free of long distance effects related the presence of four quark operators in the effective Hamiltonian (see for example [15]). It is well known that the situation for the $b \to s \ell^+ \ell^-$ is more problematic, since this decay is described by several Wilson coefficients, each bringing along its own uncertainty, and also in this decay the long distance effects are essential and should be considered. In spite of all these theoretical advantages, it might be very difficult to measure the inclusive mode $B \to X_s \nu \bar{\nu}$, because it requires to construct all X_s. Therefore it might be much easier to measure the exclusive mode $B \to K^* \nu \bar{\nu}$ experimentally. which has "large" branching ratio in SM, about 10^{-5} [16].

In this paper we study the $B \to K^* \nu \bar{\nu}$ decay for a general model independent form of effective Hamiltonian. $B \to K^* \nu \bar{\nu}$ decay has been extensively investigated in framework of the SM and its various minimal extensions [16, 17]. Note that this mode was studied in [18] in a similar way to our analysis, but in that work scalar and tensor type interactions were introduced via violation of the lepton number since neutrino was assumed to be massless. But Super Kamiokande [19] results indicated that neutrino has mass. Therefore neutrino has the right components and we can introduce scalar and tensor interactions without any lepton number violation. Another novel property of the present work is the appearence of new structures, e.g., terms that are proportional to C_{LR}^{tot} and C_{RR} (see Eq. (4) below), which are absent in [18].

The paper is organized as follows. In section 2 we give the most general, model independent form of effective Hamiltonian. We then parametrize the long distance effects by appropriate form factors and calculate the differential decay width. Section 3 is devoted to the numerical analysis and concluding remarks.

2 Differential decay width

The exclusive $B \to K^* \nu \bar{\nu}$ decay at quark level is described by $b \to s \nu \bar{\nu}$ transition. The decay amplitude for the $b \to s \nu \bar{\nu}$ decay in a general model independent form can be written as (for general form of matrix element for the $b \to s \ell^+ \ell^-$ decay, see also [20, 21])

$$
\begin{aligned}
\mathcal{M} &= \frac{G_F \alpha}{4\sqrt{2}\pi \sin^2\theta_W} V_{tb} V_{ts}^* \\
&\times \Big\{ C_{LL}^{tot}\, \bar{s}\gamma_\mu (1-\gamma_5) b \, \bar{\nu}\gamma^\mu (1-\gamma_5)\nu + C_{LR}^{tot}\, \bar{s}\gamma_\mu (1-\gamma_5) b \, \bar{\nu}\gamma^\mu (1+\gamma_5)\nu \\
&+ C_{RL}\, \bar{s}\gamma_\mu (1+\gamma_5) b \, \bar{\nu}\gamma^\mu (1-\gamma_5)\nu + C_{RR}\, \bar{s}\gamma_\mu (1+\gamma_5) b \, \bar{\nu}\gamma^\mu (1+\gamma_5)\nu \\
&+ C_{LRLR}\, \bar{s}(1+\gamma_5) b \, \bar{\nu}(1+\gamma_5)\nu + C_{RLLR}\, \bar{s}(1-\gamma_5) b \, \bar{\nu}(1+\gamma_5)\nu \\
&+ C_{LRRL}\, \bar{s}(1+\gamma_5) b \, \bar{\nu}(1-\gamma_5)\nu + C_{RLRL}\, \bar{s}(1-\gamma_5) b \, \bar{\nu}(1-\gamma_5)\nu \\
&+ C_T\, \bar{s}\sigma_{\mu\nu} b \, \bar{\nu}\sigma^{\mu\nu}\nu + i C_{TE}\, \epsilon^{\mu\nu\alpha\beta} \bar{s}\sigma_{\mu\nu} b \, \bar{\nu}\sigma_{\alpha\beta}\nu \Big\} \, . \qquad (4)
\end{aligned}
$$

Two of the four vector interactions containing C_{LL}^{tot} and C_{LR}^{tot} already exist in the SM in combinations $(C_9 - C_{10})$ and $(C_9 + C_{10})$ for the $b \to s \ell^+ \ell^-$ decay, while in the present work for the $b \to s \nu \bar{\nu}$ transition we have $C_9 = -C_{10}$. Therefore writing

$$
\begin{aligned}
C_{LL}^{tot} &= 2X + C_{LL} \, , \\
C_{LR}^{tot} &= C_{LR} \, ,
\end{aligned}
$$

one concludes that C_{LL}^{tot} describes the sum of the contributions from the SM and the new physics. The remaining coefficients C_{LRLR}, C_{LRRL}, C_{RLLR}, C_{RLRL} and C_T, C_{TE} describe the scalar and tensor interactions, respectively, which are absent in the SM. The decay amplitude (4) has a rather general form as it includes nine additional operators not found in the SM.

The decay amplitude of the semileptonic $B \to K^* \nu \bar{\nu}$ decay can be obtained after evaluating matrix elements of the quark operators in Eq. (4) between the initial $|B(p_B)\rangle$ and final $\langle K^*(p_{K^*}, \varepsilon)|$ states. It follows from Eq. (4) that we need the following matrix elements

$$\begin{aligned}
&\langle K^* \left| \bar{s}\gamma_\mu (1 \pm \gamma_5) b \right| B \rangle \ , \\
&\langle K^* \left| \bar{s}(1 \pm \gamma_5) b \right| B \rangle \ , \\
&\langle K^* \left| \bar{s}\sigma_{\mu\nu} b \right| B \rangle \ .
\end{aligned}$$

These matrix elements can be written in terms of the form factors in the following way:

$$\begin{aligned}
&\langle K^*(p_{K^*}, \varepsilon) \left| \bar{s}\gamma_\mu (1 \pm \gamma_5) b \right| B(p_B) \rangle = \\
&\quad -\epsilon_{\mu\nu\lambda\sigma} \varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma \frac{2V(q^2)}{m_B + m_{K^*}} \pm i\varepsilon_\mu^* (m_B + m_{K^*}) A_1(q^2) \qquad (5) \\
&\quad \mp i(p_B + p_{K^*})_\mu (\varepsilon^* q) \frac{A_2(q^2)}{m_B + m_{K^*}} \mp i q_\mu \frac{2m_{K^*}}{q^2} (\varepsilon^* q) \left[A_3(q^2) - A_0(q^2) \right] \ ,
\end{aligned}$$

$$\begin{aligned}
&\langle K^*(p_{K^*}, \varepsilon) \left| \bar{s}\sigma_{\mu\nu} b \right| B(p_B) \rangle = \\
&\quad i\epsilon_{\mu\nu\lambda\sigma} \Bigg\{ -2T_1(q^2) \varepsilon^{*\lambda} (p_B + p_{K^*})^\sigma + \frac{2}{q^2} (m_B^2 - m_{K^*}^2) \Big[T_1(q^2) - T_2(q^2) \Big] \varepsilon^{*\lambda} q^\sigma \qquad (6) \\
&\quad - \frac{4}{q^2} \left[T_1(q^2) - T_2(q^2) - \frac{q^2}{m_B^2 - m_{K^*}^2} T_3(q^2) \right] (\varepsilon^* q) p_{K^*}^\lambda q^\sigma \Bigg\} \ .
\end{aligned}$$

where ε is the polarization vector of K^* meson and $q = p_B - p_{K^*}$ is the momentum transfer. To ensure finiteness of (5) at $q^2 = 0$, it is usually assumed that $A_3(q^2 = 0) = A_0(q^2 = 0)$ and $T_1(q^2 = 0) = T_2(q^2 = 0)$. The matrix element $\langle K^* \left| \bar{s}(1 \pm \gamma_5) b \right| B \rangle$ can be calculated by contracting both sides of Eq. (5) with q^μ and using equation of motion. Neglecting the mass of the strange quark we get

$$\langle K^*(p_{K^*}, \varepsilon) \left| \bar{s}(1 \pm \gamma_5) b \right| B(p_B) \rangle = \frac{1}{m_b} \Big[\mp 2i m_{K^*} (\varepsilon^* q) A_0(q^2) \Big] \ . \qquad (7)$$

In deriving Eq. (7) we have used the following relation between the form factors A_1 , A_2 and A_3 (see [22])

$$A_3(q^2) = \frac{1}{2m_{K^*}} \Big[(m_B + m_{K^*}) A_1(q^2) - (m_B - m_{K^*}) A_2(q^2) \Big] \ .$$

Taking into account Eqs. (4–7), the matrix element of the $B \to K^* \bar{\nu} \nu$ decay can be written as

$$\begin{aligned}
&\mathcal{M}(B \to K^* \nu \bar{\nu}) = \frac{G_F \alpha}{4\sqrt{2}\pi \sin^2 \theta_W} V_{tb} V_{ts}^* \\
&\quad \times \Bigg\{ \bar{\nu} \gamma^\mu (1 - \gamma_5) \nu \Big[-2A \epsilon_{\mu\nu\lambda\sigma} \varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma - iB_1 \varepsilon_\mu^* + iB_2 (\varepsilon^* q)(p_B + p_{K^*})_\mu + iB_3 (\varepsilon^* q) q_\mu \Big] \\
&\quad + \bar{\nu} \gamma^\mu (1 + \gamma_5) \nu \Big[-2C \epsilon_{\mu\nu\lambda\sigma} \varepsilon^{*\nu} p_{K^*}^\lambda q^\sigma - iD_1 \varepsilon_\mu^* + iD_2 (\varepsilon^* q)(p_B + p_{K^*})_\mu + iD_3 (\varepsilon^* q) q_\mu \Big] \\
&\quad + \bar{\nu} (1 - \gamma_5) \nu \Big[iB_4 (\varepsilon^* q) \Big] + \bar{\nu} (1 + \gamma_5) \nu \Big[iB_5 (\varepsilon^* q) \Big]
\end{aligned}$$

$$+4\bar\nu\sigma^{\mu\nu}\nu\Big(iC_T\epsilon_{\mu\nu\lambda\sigma}\Big)\Big[-2T_1\varepsilon^{\ast\lambda}(p_B+p_{K^\ast})^\sigma+B_6\varepsilon^{\ast\lambda}q^\sigma-B_7(\varepsilon^\ast q){p_{K^\ast}}^\lambda q^\sigma\Big]$$
$$+16C_{TE}\bar\nu\sigma_{\mu\nu}\nu\Big[-2T_1\varepsilon^{\ast\mu}(p_B+p_{K^\ast})^\nu+B_6\varepsilon^{\ast\mu}q^\nu-B_7(\varepsilon^\ast q){p_{K^\ast}}^\mu q^\nu\Big\}~, \qquad (8)$$

where

$$\begin{aligned}
A &= (C_{LL}^{tot}+C_{RL})\frac{V}{m_B+m_{K^\ast}}~,\\
B_1 &= (C_{LL}^{tot}-C_{RL})(m_B+m_{K^\ast})A_1~,\\
B_2 &= (C_{LL}^{tot}-C_{RL})\frac{A_2}{m_B+m_{K^\ast}}~,\\
B_3 &= 2(C_{LL}^{tot}-C_{RL})m_{K^\ast}\frac{A_3-A_0}{q^2}~,\\
C &= A(C_{LL}^{tot}\to C_{LR}^{tot},~C_{RL}\to C_{RR})~,\\
D_1 &= B_1(C_{LL}^{tot}\to C_{LR}^{tot},~C_{RL}\to C_{RR})~,\\
D_2 &= B_2(C_{LL}^{tot}\to C_{LR}^{tot},~C_{RL}\to C_{RR})~,\\
D_3 &= B_3(C_{LL}^{tot}\to C_{LR}^{tot},~C_{RL}\to C_{RR})~,\\
B_4 &= -2(C_{LRRL}-C_{RLRL})\frac{m_{K^\ast}}{m_b}A_0~,\\
B_5 &= -2(C_{LRLR}-C_{RLLR})\frac{m_{K^\ast}}{m_b}A_0~,\\
B_6 &= 2(m_B^2-m_{K^\ast}^2)\frac{T_1-T_2}{q^2}~,\\
B_7 &= \frac{4}{q^2}\left(T_1-T_2-\frac{q^2}{m_B^2-m_{K^\ast}^2}T_3\right)~.
\end{aligned}$$

Note that in further calculations we set neutrino mass to zero.

From the matrix element (8) it is straightforward to derive the missing mass–squared spectrum corresponding to the longitudinally and transversally polarized $K^\ast$ meson. In the case of longitudinally polarized $K^\ast$ meson, we get for the missing mass–squared spectrum

$$\begin{aligned}
\frac{d\Gamma_L}{dq^2} &= N_\nu\left[\frac{G_F\alpha}{4\sqrt{2}\pi\sin^2\theta_W}\right]^2\left|V_{tb}V_{ts}^\ast\right|^2\frac{1}{256m_B^2\pi^3}\,\frac{1}{3m_{K^\ast}^2}\\
&\times \Big\{4\left|2B_1h+B_2\lambda\right|^2+4\left|2D_1h+D_2\lambda\right|^2+6\Big(\left|B_4\right|^2+\left|B_5\right|^2\Big)\lambda q^2\\
&+ 16\left|4B_6h-B_7\lambda\right|^2\Big(4\left|C_{TE}\right|^2+\left|C_T\right|^2\Big)q^2\times\Big[16\left|T_1\right|^2(m_B^2+3m_{K^\ast}^2-q^2)^2\\
&- 16\mathrm{Re}(B_6T_1^\ast)\Big[\lambda+4m_{K^\ast}^2(m_B^2-m_{K^\ast}^2)\Big]+8\mathrm{Re}(B_7T_1^\ast)(\lambda+3m_{K^\ast}^2-q^2)\Big]\Big\}~. \qquad (9)
\end{aligned}$$

For the transversally polarized $K^\ast$ meson, the differential decay width takes the following form

$$\frac{d\Gamma_\mp}{dq^2} = N_\nu\left[\frac{G_F\alpha}{4\sqrt{2}\pi\sin^2\theta_W}\right]^2\left|V_{tb}V_{ts}^\ast\right|^2$$

$$
\begin{aligned}
\times & \quad \left\{ \frac{16}{3}\sqrt{\lambda}\, q^2 \Big(\left|A \pm B_1\right|^2 + \left|C \pm D_1\right|^2 \Big) \right. \\
+ & \quad \frac{2048}{3}\sqrt{\lambda}\,\mathrm{Re}(C_T C_{TE}^*)\Big[\mp 4(m_B^2 - m_{K^*}^2)\left|T_1\right|^2 + 2\mathrm{Re}(B_6 T_1^*) q^2 \Big] \\
+ & \quad \frac{256}{3}\Big(4\left|C_{TE}\right|^2 + \left|C_T\right|^2 \Big)\Big(\left|B_6\right|^2 q^4 + 4[\lambda + (m_B^2 - m_{K^*}^2)]\left|T_1\right|^2 \\
- & \quad \left. 4(m_B^2 - m_{K^*}^2)\mathrm{Re}(B_6 T_1^*) q^2 \Big) \right\} \qquad (10)
\end{aligned}
$$

In Eqs. (9) and (10) $N_\nu = 3$ is the number of light neutrinos,

$$
\begin{aligned}
\lambda(m_B^2, m_{K^*}^2, q^2) &= m_B^4 + m_{K^*}^4 + q^4 - 2m_B^2 q^2 - 2m_{K^*}^2 q^2 - 2m_B^2 m_{K^*}^2 \;, \\
h &= \frac{1}{2}(m_B^2 - m_{K^*}^2 - q^2) \;.
\end{aligned}
$$

It should be noted that in experiments due to the non–detectibility of the neutrinos, it is impossible to discriminate the transverse polarization +1 from −1. For this reason these two polarization states must be added, i.e.,

$$
\frac{d\Gamma_T}{dq^2}(B \to K^* \nu\bar{\nu}) = \frac{d\Gamma_+}{dq^2}(B \to K^* \nu\bar{\nu}) + \frac{d\Gamma_-}{dq^2}(B \to K^* \nu\bar{\nu}) \;. \qquad (11)
$$

From Eqs. (9) and (10) we observe that scalar interaction gives contribution to the differential decay width $d\Gamma_L/dq^2$ when K^* is longitudinally polarized and does not contribute to $d\Gamma_\pm/dq^2$ when it transversally polarized. This result can be explained in the following way. When B meson is at rest, K^* meson and neutrino pair must be in flight along opposite directions. When K^* meson is transversally polarized the total helicity of the neutrino pair must be ± 1, since the initial B meson spin is equal to zero. But from the neutrino antineutrino pair, which are in flight along the same direction, one can organize a total helicity of ± 1 by flipping one of the neutrino's helicity. But this spin flip can be achieved by inserting mass of neutrino. However, as has already been mentioned previously, we neglect the neutrino mass in our calculations and for this reason in the expression for the differential decay width when K^* meson is transversally polarized, the terms describing the scalar interaction disappear.

3 Numerical analysis

In this section we will study the sensitivity of the branching ratio and missing mass–squared spectrum to the new Wilson coefficients.

In performing numerical calculations, as can easily be seen from Eqs. (9) and (10), first of all, we need the expressions for the form factors. For the values of the form factors, we have used the results of [23] (see also [24] and [25]), where the radiative corrections to the leading twist contribution and $SU(3)$ breaking effects are also taken into account. It is shown in [23] that the q^2 dependence of the form factors can be represented in terms of

three parameters as

$$F(q^2) = \frac{F(0)}{1 - a_F \dfrac{q^2}{m_B^2} + b_F \left(\dfrac{q^2}{m_B^2} \right)^2} \ ,$$

where, the values of parameters $F(0)$, a_F and b_F for the $B \to K^*$ decay are listed in Table 1.

	$F(0)$	a_F	b_F
$A_1^{B\to K^*}$	0.34 ± 0.05	0.60	-0.023
$A_2^{B\to K^*}$	0.28 ± 0.04	1.18	0.281
$V^{B\to K^*}$	0.46 ± 0.07	1.55	0.575
$T_1^{B\to K^*}$	0.19 ± 0.03	1.59	0.615
$T_2^{B\to K^*}$	0.19 ± 0.03	0.49	-0.241
$T_3^{B\to K^*}$	0.13 ± 0.02	1.20	0.098

Table 1: B meson decay form factors in a three-parameter fit, where the radiative corrections to the leading twist contribution and SU(3) breaking effects are taken into account [25].

In Figs. (1) and (2) we present the dependence of the branching ratios on the new Wilson coefficients, when K^* polarized transversally and longitudinally, respectively. From both figures we see that when C_{LL} lies in the region from -4 to 0 (in numerical calculations all new Wilson coefficients vary in the range from -4 to $+4$), branching ratios are lower than the SM prediction. Moreover, when C_{LL} increases from 0 up to $+4$, branching ratios become larger than the SM result and for increasing values of C_{LL} the departure from SM becomes substantial. This behavior can be explained by the fact that in the range from -4 to 0 the new Wilson coefficient C_{LL} gives destructive and in the second half of the range from 0 to $+4$ it gives constructive interference to the SM result. For the Wilson coefficients C_{RR} and C_{LR}, we observe the following dependence of the branching ratios. Up to the zero value of the Wilson coefficients the branching ratios decrease and at $C_{LR} = C_{RR} = 0$ they coincide with the SM prediction. Furthermore, with increasing C_{RR}, C_{LR} both $\mathcal{B}_L$ and $\mathcal{B}_T$ increase. Qualitatively, this behavior could be explained as follows. When all Wilson coefficients, except C_{RR} (or C_{LR}), are zero, $\mathcal{B}_L$ and $\mathcal{B}_T$ are proportional to $|C_{RR}|^2$ (or $|C_{LR}|^2$). Therefore as C_{RR} (or C_{LR}) increases in the region from -4 to 0, $\mathcal{B}_L$ and $\mathcal{B}_T$ decrease and when C_{RR} (or C_{LR}) increase from 0 to $+4$ the above–mentioned grow larger. Obviously the dependence of $\mathcal{B}_L$ and $\mathcal{B}_T$ on C_{RR} (or C_{LR}) must be symmetric, and this expectation is confirmed by the numerical calculations. In the case of the dependence of branching ratios on the Wilson coefficient C_{RL}, we observe that $\mathcal{B}_L$ decreases with changing values of C_{RL} in the range from -4 to $+4$. We can argue about this dependence as follows. For this Wilson coefficient $\mathcal{B}_L$ is proportional to $|2X - C_{RL}|^2$, and hence $\mathcal{B}_L$ decreases for

increasing values of C_{RL}, as expected. However the situation is different for the $\mathcal{B}_T$, as can easily be seen from the figure, it decreases when C_{RL} increases from -4 to 0 and then increases in the positive half of the range.

As has already been noted, $\mathcal{B}_T$ is independent of the scalar type interaction while $\mathcal{B}_L$ is dependent. From Fig. (2) we observe that for all scalar type interaction coefficients the branching ratio $\mathcal{B}_L$ shows a similar behavior, i.e., it decreases in the first half of the range of variation of the scalar interaction coefficients and increases for the positive part of the range from 0 to +4. In contrary to the previous cases, as Figs. (1) and (2) depict, $\mathcal{B}_L$ and $\mathcal{B}_T$ show quite a strong dependence on the tensor type interaction coefficients C_T and C_{TE}. As can easily be seen from Eqs. (9) and (10), $\mathcal{B}_L$ and $\mathcal{B}_T$ depend as moduli square on C_T and C_{TE} and therefore this dependence must be symmetric. If we assume that the departure from SM prediction is expected to be small, it will put very strong restriction to the tensor type interactions.

We also analyze the missing mass–squared spectrum on new Wilson coefficients. All qualitative arguments which we have put forward in discussing the dependence of branching ratios on new Wilson coefficients work their way similarly and remains in power in the case of missing mass–squared spectrum as well. As an example in Fig. (3) we present the dependence of missing mass–squared spectrum at four different values of C_{LL}, namely $-2,~-1~,0,~+1,~+2$.

Finally a few words about the dependence of the another experimentally measurable quantity, namely $\mathcal{B}_L/\mathcal{B}_T$, on the new Wilson coefficients are in order. Our numerical analysis shows that this ratio is practically independent of the new Wilson coefficients and are very close to the SM prediction. Therefore study of this ratio can not serve as an effective tool in search of new physics beyond SM.

In conclusion, using the most general, model independent form of the effective Hamiltonian we have studied the sensitivity of the branching ratios $\mathcal{B}_L$ and $\mathcal{B}_T$ to the new Wilson coefficients. The main result of this study is that the branching ratios and the missing mass–squared spectrum are very useful in looking new physics beyond SM.

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Figure captions

Fig. (1) The dependence of the branching ratio of the $B \to K^* \nu \bar{\nu}$ decay on the new Wilson coefficients, when K^* polarized transversally. The line indicated by C_{XXXX} denotes any one of the four scalar interaction coefficients, namely, C_{LRRL}, C_{RLLR}, C_{LRLR} and C_{RLRL}.

Fig. (2) The same as in Fig. (1), but when K^* polarized longitudinally.

Fig. (3) The dependence of missing mass–squared spectrum at four different values of C_{LL}, namely -2, -1 ,0, $+1$, $+2$. The first four lines represent the case when K^* polarized longitudinally, and the remaining four lines represent the case when K^* polarized transversally.

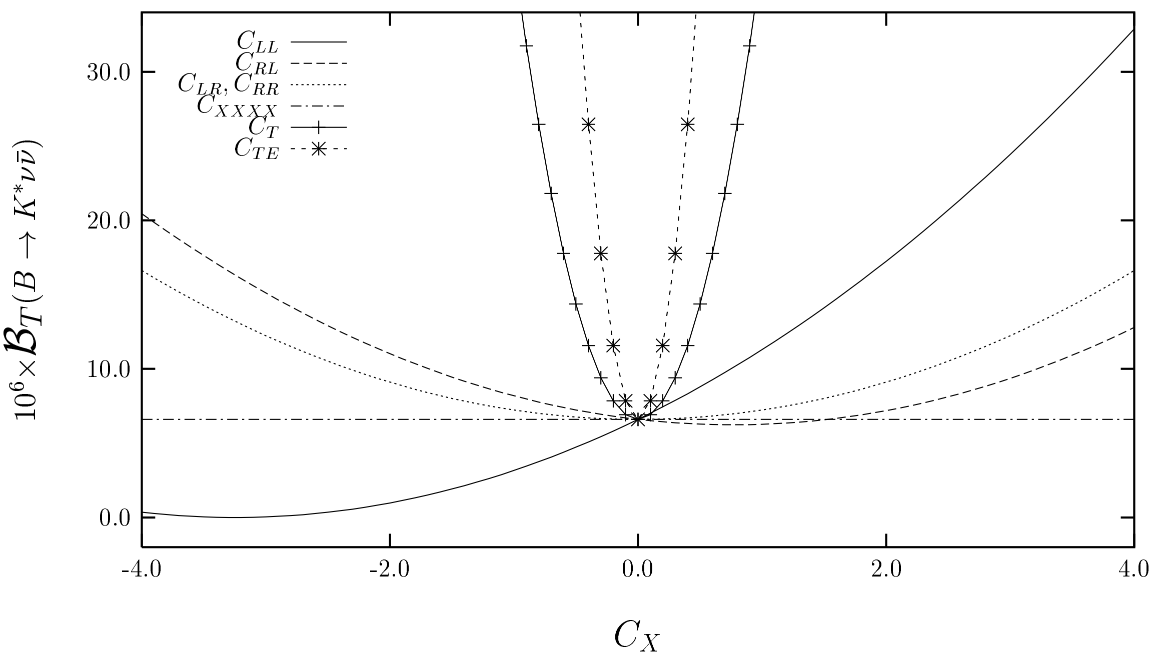


Figure 1:

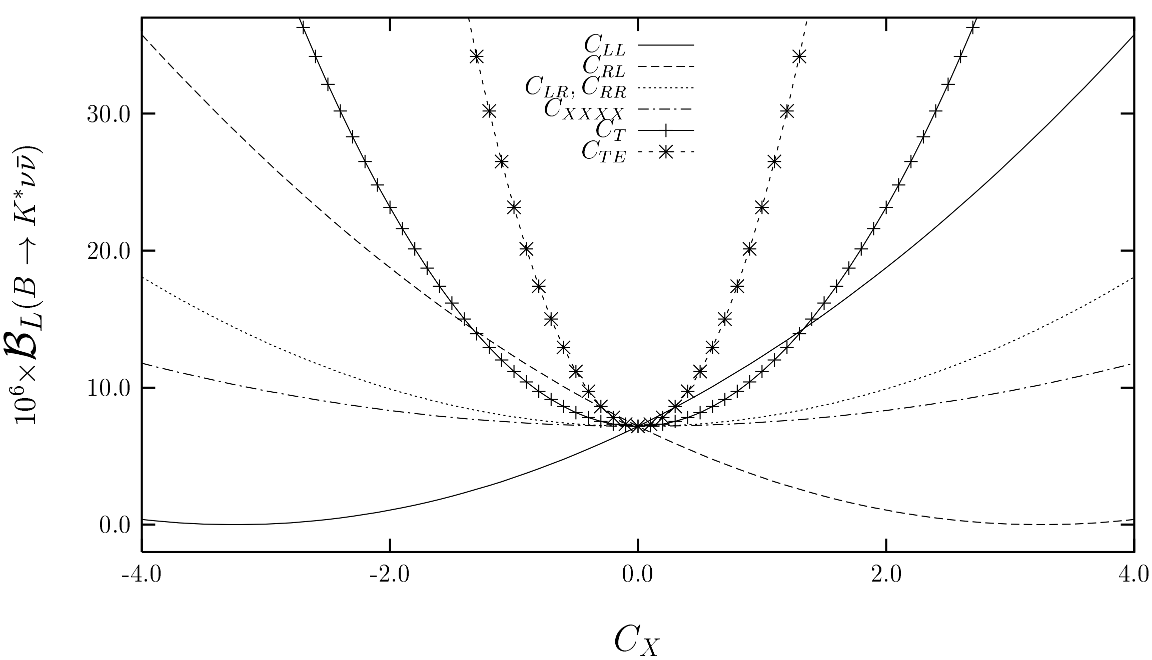
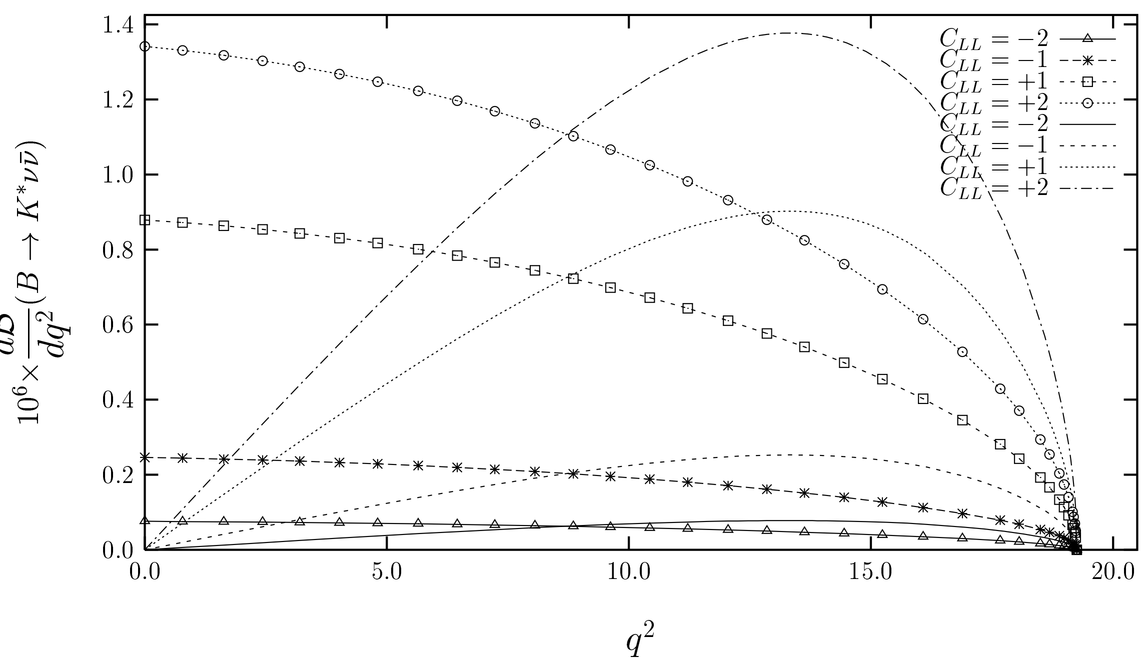


Figure 2:



$C_{LL}=-2$
$C_{LL}=-1$
$C_{LL}=+1$
$C_{LL}=+2$
$C_{LL}=-2$
$C_{LL}=-1$
$C_{LL}=+1$
$C_{LL}=+2$
1.4
1.2
1.0
0.8
0.6
0.4
0.2
0.0
0.0
5.0
10.0
15.0
20.0
$10^6\times\frac{d\mathcal{B}}{dq^2}(B\to K^*\nu\bar{\nu})$
q^2

Figure 3: